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Weak decays $\Xi_Q \rightarrow \Lambda_Q \pi$

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Abstract

The weak decays $\Xi_b \rightarrow \Lambda_b \pi$ and $\Xi_c \rightarrow \Lambda_c \pi$, in which the heavy quark is not destroyed, are discussed. The branching fractions for these decays, corresponding to an absolute rate of order 0.01 ps^{-1} , should be at a one percent level for the b hyperons and at a (few) per mill level for Ξ_c , possibly making feasible their experimental study in future. It is shown, through an application of the heavy quark limit, the flavor SU(3) symmetry, and PCAC, that the $\Delta I = 1/2$ rule should hold very well in these decays, and also that the Ξ_b decays are purely S wave in the symmetry limit, while the difference between the S wave amplitudes of the Ξ_c decays and those for the Ξ_b is related, in terms of the heavy quark expansion, to the difference of the total decay rates within the (Ξ_c, Λ_c) triplet of charmed hyperons. We also comment on the amplitudes of the semileptonic transitions $\Xi_Q \rightarrow \Lambda_Q \ell \nu$ and on the weak radiative decays $\Xi_Q \rightarrow \Lambda_Q \gamma$.

The absolutely dominant processes in weak decays of the b and c hadrons are, naturally, those associated with the decay of the heavy quark. It is with these decays where, understandably, lies the main interest of the current phenomenological studies, with the prospects for precision determination of the CKM mixing matrix and for uncovering a CP violation in B mesons. The present paper however deals with a rather sub-dominant type of decay of strange heavy hyperons into non-strange ones, a process closely analogous to the decays of ordinary strange hyperons, and associated with the decay of the strange, rather than heavy, quark. These decays are interesting for at least two reasons: one is that the branching fractions of the decays $\Xi_b \rightarrow \Lambda_b \pi$ and $\Xi_c \rightarrow \Lambda_c \pi$ are not hopelessly small and one may expect that these decays can be studied experimentally, inasmuch as it will be feasible to study any exclusive nonleptonic decays of the b and c cascade hyperons¹, and the other reason being that these decays provide a case for a study of the ‘old’ physics in a new setting, namely a study of the structure of baryons containing one heavy quark. Thus these decays offer a testing ground for a combination of the ‘older’ methods, such as the flavor SU(3) symmetry and the current algebra with the ‘newer’ theoretical ideas related to the heavy quark limit. Moreover, as will be shown, the difference between the amplitudes of the decays $\Xi_c \rightarrow \Lambda_c \pi$ and $\Xi_b \rightarrow \Lambda_b \pi$ is related, through PCAC and the SU(3) symmetry, to the matrix elements of four-quark operators, that govern, within the heavy quark expansion, the differences of the total inclusive weak decay rates within the triplets of the heavy baryons: (Ξ_c, Λ_c) and (Ξ_b, Λ_b) . The latter matrix elements are a crucial ingredient in understanding the pre-asymptotic effects in the inclusive decay rates of heavy baryons, which are discussed with a recently renewed interest in relation to the data [1] on $\tau(\Lambda_b)/\tau(B^0)$ (for a most recent mini-review see Ref.[2], see also the recent papers [3, 4]).

A very approximate estimate of the absolute rates of the discussed decays can be done by comparing them to similar strangeness-changing decays of ordinary strange hyperons, with rates typically of order $0.01ps^{-1}$. For the charmed hyperons the mass difference between Ξ_c and Λ_c is known [1] to be about 180 MeV, i.e. quite close to the mass differences of ordinary hyperons differing by one unit of strangeness, and the mass difference between Ξ_b and Λ_b should be very close to that for the charmed hyperons due to the heavy quark limit considerations:

$$M(\Xi_b) - M(\Lambda_b) = M(\Xi_c) - M(\Lambda_c) + O(m_c^{-2} - m_b^{-2}), \quad (1)$$

¹The branching ratio $B(\Xi_b^- \rightarrow \Lambda_b \pi^-)$ may well exceed 1%, thus possibly being the largest among the branching fractions for individual exclusive nonleptonic decay channels of the Ξ_b^- .

since there are no terms of order $1/m_Q$ in this mass splitting. In a more detailed consideration, to be discussed in this paper, the baryonic matrix elements, determining the decay amplitudes, are somewhat different from those for decays of ordinary hyperon decays, thus the specific absolute rates can differ from the simplistic estimate. However, using the mentioned relation of the difference of the amplitudes for Ξ_c and Ξ_b decays to the difference of the lifetimes of the two Ξ_c hyperons and the Λ_c , we find that this difference alone would correspond to the rate of, e.g. the decay $\Xi_c^0 \rightarrow \Lambda_c \pi^-$, equal to $9 \times 10^{-3} ps^{-1}$, i.e. in the same range as the simplistic estimate. Thus comparing these estimates for the absolute rates with the known lifetimes of the Ξ_c and Ξ_b hyperons, one concludes that the discussed decays should have branching fractions at a per mill level for the Ξ_c hyperons and at a one percent level for the Ξ_b .

In addition to the strangeness changing decays of heavy cascade hyperons with emission of a pion, we consider here, as a ‘by-product’, two other types of $\Xi_Q \rightarrow \Lambda_Q$ transitions: with emission of a lepton pair and with emission of a photon, which both of course have much lower probability than the pion transitions. For the semileptonic transitions we note that the axial form factor of the weak current is zero in the heavy quark limit for both $\Xi_b \rightarrow \Lambda_b \ell \nu$ and $\Xi_c \rightarrow \Lambda_c \ell \nu$, while the vector form factor is $g_V = 1$ in the flavor SU(3) limit (with the usual further consequences of the Ademollo-Gatto theorem for the SU(3) violation effects). For the photon transitions it is concluded here that the decay $\Xi_b \rightarrow \Lambda_b \gamma$ is forbidden in the heavy quark limit, while this generally is not the case for $\Xi_c \rightarrow \Lambda_c \gamma$.

The discussed transitions among the heavy hyperons are induced by two underlying weak processes: the ‘spectator’ decay of a strange quark, $s \rightarrow u \bar{u} d$, $s \rightarrow u \ell \nu$, or $s \rightarrow d \gamma$, which does not involve the heavy quark, and the ‘non-spectator’ weak scattering (WS)

$$s c \rightarrow c d \tag{2}$$

through the weak interaction of the $c \rightarrow d$ and $s \rightarrow c$ currents. One can also readily see that the WS mechanism contributes only to the decays $\Xi_c \rightarrow \Lambda_c \pi$ and generally, through a photon emission in WS, to the radiative transition $\Xi_c^+ \rightarrow \Lambda_c \gamma$, while the semileptonic decay $\Xi_c^0 \rightarrow \Lambda_c \ell \nu$ and all the transitions between the Ξ_b hyperons and the Λ_b are contributed only by the ‘spectator’ processes.

An important starting point in considering the transitions $\Xi_Q \rightarrow \Lambda_Q$ induced by the ‘spectator’ decay of the strange quark, is that in the heavy quark limit the spin of the heavy quark completely decouples from the spin variables of the light component of the baryon, and that the latter light component in both the initial and the final baryon forms a $J^P = 0^+$

state with the quantum numbers of a diquark. Thus these transitions are of a $0^+ \rightarrow 0^+$ type, which imposes strong constraints on the decay amplitudes. In particular, for the pion emission, these constraints imply that the decay amplitude is purely S wave, while the P wave amplitude is zero in the limit of infinite heavy quark mass². The implication for the ‘spectator’ radiative transition is that it is forbidden in this limit, thus predicting a strong suppression of the decay $\Xi_b^0 \rightarrow \Lambda_b \gamma$. For the semileptonic decays, contributed only by the ‘spectator’ decay, the constraint from the $0^+ \rightarrow 0^+$ transition is that the axial hadronic form factor is vanishing, $g_A = 0$, while the vector form factor is $g_V = 1$. The corresponding decay rate $\Gamma(\Xi_Q \rightarrow \Lambda_Q e \nu) = G_F^2 \sin^2 \theta_c (\Delta M)^5 / (60\pi^3) \approx 1.0 \times 10^6 s^{-1}$ is however too small to be of a possible phenomenological relevance in the nearest future.

For further consideration of the pion transitions $\Xi_Q \rightarrow \Lambda_Q \pi$ we write the well known expression for the nonleptonic strangeness-changing weak Hamiltonian (see e.g. [5])

$$H_W = \sqrt{2} G_F \sin \theta_c \left\{ (C_+ + C_-) \left[(\bar{u}_L \gamma_\mu s_L) (\bar{d}_L \gamma_\mu u_L) - (\bar{c}_L \gamma_\mu s_L) (\bar{d}_L \gamma_\mu c_L) \right] + (C_+ - C_-) \left[(\bar{d}_L \gamma_\mu s_L) (\bar{u}_L \gamma_\mu u_L) - (\bar{d}_L \gamma_\mu s_L) (\bar{c}_L \gamma_\mu c_L) \right] \right\} . \quad (3)$$

In this formula the weak Hamiltonian is assumed to be normalized (in LLO) at $\mu = m_c$, so that the renormalization coefficients are $C_- = C_+^{-2} = (\alpha_s(m_c)/\alpha_s(m_W))^{12/25}$. The terms in the Hamiltonian (3) without the charmed quark fields describe the ‘spectator’ nonleptonic decay of the strange quark, while those with the c quark correspond to the WS process (2). It should be noted that the part of H_W with (virtual) charmed quarks indirectly contributes to the ‘spectator’ process as well by providing a GIM cutoff for the ‘penguin’ mechanism [6] at $\mu = m_c$. However at any normalization point μ below m_c this part does not explicitly show up, and for this reason we refer to the terms of H_W without the charmed quark fields as ‘spectator’ ones and those with c and \bar{c} as ‘non-spectator’ ones. One could evolve the weak Hamiltonian down to a low normalization point μ , such that $\mu \ll m_c$ to make this separation explicit (the ‘spectator’ part then evolves according to the treatment in Ref.[6], while the evolution of the ‘non-spectator’ part is described by the ‘hybrid’ anomalous dimension [7], and is essentially equivalent to that considered in [8]). However, the present paper makes no attempt at constructing models for the hadronic matrix elements at a low μ , thus writing the corresponding formulas here would be redundant and it is quite sufficient for our purposes to use the expression (3) at $\mu = m_c$ with the separation between the ‘spectator’ and the ‘non-spectator’ parts as noted.

²For a general phenomenological treatment of the amplitudes of hyperon pion transitions, $B' \rightarrow B \pi$, in terms of partial waves, see e.g. the textbook [5].

As discussed above, the ‘spectator’ process gives rise only to the S wave amplitudes of the decays $\Xi_Q \rightarrow \Lambda_Q \pi$, while the ‘non-spectator’ part involves the spin of the charmed quark, and generally may induce a P wave as well as an S wave in the decays of the Ξ_c hyperons. According to the well known current algebra technique, the S wave amplitudes of pion emission can be considered in the chiral limit at zero four-momentum of the pion, where they are described by the PCAC reduction formula (pole terms are absent in these processes):

$$\langle \Lambda_Q \pi_i(p=0) | H_W | \Xi_Q \rangle = \frac{\sqrt{2}}{f_\pi} \langle \Lambda_Q | [Q_i^5, H_W] | \Xi_Q \rangle, \quad (4)$$

where π_i is the pion triplet in the Cartesian notation, and Q_i^5 is the corresponding isotopic triplet of axial charges. The constant $f_\pi \approx 130 \text{ MeV}$, normalized by the charged pion decay, is used here, hence the coefficient $\sqrt{2}$ in eq.(4).

It is straightforward to see from eq.(4) that in the PCAC limit the discussed decays should obey the $\Delta I = 1/2$ rule. Indeed, the commutator of the weak Hamiltonian with the axial charges transforms under the isotopic SU(2) in the same way as the Hamiltonian itself. In other words, the $\Delta I = 1/2$ part of H_W after the commutation gives an $\Delta I = 1/2$ operator, while the $\Delta I = 3/2$ part after the commutation gives an $\Delta I = 3/2$ operator. The latter operator however cannot have a non vanishing matrix element between an isotopic singlet, Λ_Q , and an isotopic doublet, Ξ_Q . Thus the $\Delta I = 3/2$ part of H_W gives no contribution to the S wave amplitudes in the PCAC limit. The P wave part of the amplitude vanishes at a zero pion momentum and thus is not given by eq.(4). However, as discussed, the P wave can arise only in the decays of charmed hyperons and only from the ‘non-spectator’ part of the H_W , which is pure $\Delta I = 1/2$ to start with.

Once the isotopic properties of the decay amplitudes are fixed, one can concentrate on specific charge decay channels, e.g. $\Xi_b^- \rightarrow \Lambda_b \pi^-$ and $\Xi_c^0 \rightarrow \Lambda_c \pi^-$. An application of the PCAC relation (4) with the Hamiltonian from eq.(3) to these decays, gives the expressions for the amplitudes at $p=0$ in terms of baryonic matrix elements of four-quark operators:

$$\begin{aligned} & \langle \Lambda_b \pi^-(p=0) | H_W | \Xi_b^- \rangle = \\ & \frac{\sqrt{2}}{f_\pi} G_F \sin \theta_c \langle \Lambda_b | (C_+ + C_-) \left[(\bar{u}_L \gamma_\mu s_L) (\bar{d}_L \gamma_\mu d_L) - (\bar{u}_L \gamma_\mu s_L) (\bar{u}_L \gamma_\mu u_L) \right] + \\ & (C_+ - C_-) \left[(\bar{d}_L \gamma_\mu s_L) (\bar{u}_L \gamma_\mu d_L) - (\bar{u}_L \gamma_\mu s_L) (\bar{u}_L \gamma_\mu u_L) \right] | \Xi_b^- \rangle = \\ & \frac{\sqrt{2}}{f_\pi} G_F \sin \theta_c \langle \Lambda_b | C_- \left[(\bar{u}_L \gamma_\mu s_L) (\bar{d}_L \gamma_\mu d_L) - (\bar{d}_L \gamma_\mu s_L) (\bar{u}_L \gamma_\mu d_L) \right] - \\ & \frac{C_+}{3} \left[(\bar{u}_L \gamma_\mu s_L) (\bar{d}_L \gamma_\mu d_L) + (\bar{d}_L \gamma_\mu s_L) (\bar{u}_L \gamma_\mu d_L) + 2 (\bar{u}_L \gamma_\mu s_L) (\bar{u}_L \gamma_\mu u_L) \right] | \Xi_b^- \rangle, \quad (5) \end{aligned}$$

where in the last transition the operator structure with $\Delta I = 3/2$ giving a vanishing contribution is removed and only the structures with explicitly $\Delta I = 1/2$ are retained, and

$$\begin{aligned} \langle \Lambda_c \pi^- (p=0) | H_W | \Xi_c^0 \rangle &= \langle \Lambda_b \pi^- (p=0) | H_W | \Xi_b^- \rangle + \\ &\frac{\sqrt{2}}{f_\pi} G_F \sin \theta_c \langle \Lambda_c | (C_+ + C_-) (\bar{c}_L \gamma_\mu s_L) (\bar{u}_L \gamma_\mu c_L) + \\ &(C_+ - C_-) (\bar{u}_L \gamma_\mu s_L) (\bar{c}_L \gamma_\mu c_L) | \Xi_c^0 \rangle . \end{aligned} \quad (6)$$

In the latter formula the first term on the r.h.s. expresses the fact that in the heavy quark limit the ‘spectator’ amplitudes do not depend on the flavor or the mass of the heavy quark³. The rest of the expression (6) describes the ‘non-spectator’ contribution to the S wave of the charmed hyperon decay. Using the flavor SU(3) symmetry this contribution can be related to the difference of lifetimes within the charmed hyperon triplet as follows.

Due to the absence of correlation of the spin of the heavy quark in the hyperons with its light ‘environment’, the terms, involving the axial current of the charmed quark in the operators in eq.(6), give no contribution to the matrix elements. Thus the only relevant matrix elements are

$$\langle \Lambda_c | (\bar{c} \gamma_\mu c) (\bar{u} \gamma_\mu s) | \Xi_c^0 \rangle = -x \quad \text{and} \quad \langle \Lambda_c | (\bar{c}_i \gamma_\mu c_k) (\bar{u}_k \gamma_\mu s_i) | \Xi_c^0 \rangle = -y , \quad (7)$$

where, by the SU(3) symmetry, the quantities x and y coincide with those introduced in [9] in terms of differences of diagonal matrix elements over the hyperons:

$$\begin{aligned} x &= \left\langle \frac{1}{2} (\bar{c} \gamma_\mu c) [(\bar{u} \gamma_\mu u) - (\bar{s} \gamma_\mu s)] \right\rangle_{\Xi_c^0 - \Lambda_c} = \left\langle \frac{1}{2} (\bar{c} \gamma_\mu c) [(\bar{s} \gamma_\mu s) - (\bar{d} \gamma_\mu d)] \right\rangle_{\Lambda_c - \Xi_c^+} , \quad (8) \\ y &= \left\langle \frac{1}{2} (\bar{c}_i \gamma_\mu c_k) [(\bar{u}_k \gamma_\mu u_i) - (\bar{s}_k \gamma_\mu s_i)] \right\rangle_{\Xi_c^0 - \Lambda_c} = \left\langle \frac{1}{2} (\bar{c}_i \gamma_\mu c_k) [(\bar{s}_k \gamma_\mu s_i) - (\bar{d}_k \gamma_\mu d_i)] \right\rangle_{\Lambda_c - \Xi_c^+} \end{aligned}$$

with the notation for the differences of the matrix elements: $\langle \mathcal{O} \rangle_{A-B} = \langle A | \mathcal{O} | A \rangle - \langle B | \mathcal{O} | B \rangle$.

In terms of the quantities x and y in eq.(7) the difference of the S wave decay amplitudes from eq.(6) is written as

$$\begin{aligned} \Delta A_S \equiv \langle \Lambda_c \pi^- (p=0) | H_W | \Xi_c^0 \rangle - \langle \Lambda_b \pi^- (p=0) | H_W | \Xi_b^- \rangle &= \\ &\frac{G_F \sin \theta_c}{2 \sqrt{2} f_\pi} [(C_- - C_+) x - (C_+ + C_-) y] . \end{aligned} \quad (9)$$

On the other hand, the same quantities defined by eqs.(8) describe within the heavy quark expansion [8] the differences of the inclusive weak decay rates within the triplet of the

³The non-relativistic normalization for the heavy quark field is used here, corresponding to $\langle Q | Q^\dagger Q | Q \rangle = 1$. Thus the amplitudes do not contain normalization factors related to the heavy quark mass.

charmed baryons [4]. These quantities were extracted [4] from the current data on the lifetime differences for the charmed baryons. In particular it was found that the naive quark model relation $x = -y$ between the (μ independent) x and the (μ dependent) y does not hold at any μ below m_c . The numerical value of x is found as $x = -(0.04 \pm 0.01) GeV^3$, while the value of y at $\mu = m_c$ is found to be $y = 0.019 \pm 0.009 GeV^3$, which values can be directly used in eq.(9). Because of a correlation in the errors in these estimates and for possible future reference to (hopefully) more precise future data on the lifetimes, it is rather appropriate to express the difference of the amplitudes ΔA_S in eq.(9) directly in terms of the lifetimes of the charmed hyperons, using the formulas from Ref.[4]. In terms of the operators normalized at $\mu = m_c$ the relations for the differences of the total decay rates, including the dominant Cabibbo-unsuppressed nonleptonic decays as well as the decays with single Cabibbo suppression and the semileptonic decays, read as

$$\begin{aligned}\Gamma(\Xi_c^0) - \Gamma(\Lambda_c) &= \frac{G_F^2 m_c^2}{4\pi} \cos^2 \theta_c \left\{ -x \left[\cos^2 \theta_c C_+ C_- + \frac{\sin^2 \theta_c}{4} (6 C_+ C_- + 5 C_+^2 + 5 C_-^2) \right] + \right. \\ &y \left[3 \cos^2 \theta_c C_+ C_- + \frac{3 \sin^2 \theta_c}{4} (6 C_+ C_- - 3 C_+^2 + C_-^2) + 2 \right] \left. \right\} \approx \frac{G_F^2 m_c^2}{4\pi} (-1.39 x + 5.56 y) , \\ \Gamma(\Lambda_c) - \Gamma(\Xi_c^+) &= \frac{G_F^2 m_c^2}{4\pi} \left\{ -x \frac{\cos^4 \theta_c}{4} (5 C_+^2 + 5 C_-^2 - 2 C_+ C_-) + \right. \\ &y \left[\frac{3 \cos^4 \theta_c}{4} (C_-^2 - 3 C_+^2 - 2 C_+ C_-) - 2 (\cos^2 \theta_c - \sin^2 \theta_c) \right] \left. \right\} \approx -\frac{G_F^2 m_c^2}{4\pi} (2.88 x + 3.16 y) ,\end{aligned}\tag{10}$$

where for simplification of the subsequent relation the numerical values are substituted for the coefficients C_+ and C_- : $C_+ = 0.80$, $C_- = 1.55$, corresponding to $\alpha_s(m_c)/\alpha_s(m_W) = 2.5$.

The relations (10) allow to eliminate the quantities x and y from the expression (9) in favor of the differences of the total decay rates:

$$\begin{aligned}\Delta A_S &\approx -\frac{\sqrt{2} \pi \sin \theta_c}{G_F m_c^2 f_\pi} \left[0.45 (\Gamma(\Xi_c^0) - \Gamma(\Lambda_c)) + 0.04 (\Gamma(\Lambda_c) - \Gamma(\Xi_c^+)) \right] = \\ &-10^{-7} \left[0.97 (\Gamma(\Xi_c^0) - \Gamma(\Lambda_c)) + 0.09 (\Gamma(\Lambda_c) - \Gamma(\Xi_c^+)) \right] \left(\frac{1.4 GeV}{m_c} \right)^2 ps ,\end{aligned}\tag{11}$$

where, clearly, in the latter form the widths are assumed to be expressed in ps^{-1} , and $m_c = 1.4 GeV$ is used as a ‘reference’ value for the charmed quark mass. It is seen from eq.(11) that the evaluation of the difference of the amplitudes within the discussed approach is mostly sensitive to the difference of the decay rates of Ξ_c^0 and Λ_c , with only very little sensitivity to the total decay width of Ξ_c^+ . Using the current data [1] on the total decay

rates: $\Gamma(\Lambda_c) = 4.85 \pm 0.28 \text{ ps}^{-1}$, $\Gamma(\Xi_c^0) = 10.2 \pm 2 \text{ ps}^{-1}$, and the updated value [10] $\Gamma(\Xi_c^+) = 3.0 \pm 0.45 \text{ ps}^{-1}$, the difference ΔA_S is estimated as

$$\Delta A_S = -(5.4 \pm 2) \times 10^{-7} , \quad (12)$$

with the uncertainty being dominated by the experimental error in the lifetime of Ξ_c^0 . An S wave amplitude A_S of the magnitude, given by the central value in eq.(12) would produce a decay rate $\Gamma(\Xi_Q \rightarrow \Lambda_Q \pi) = |A_S|^2 p_\pi / (2\pi) \approx 0.9 \times 10^{10} \text{ s}^{-1}$, which result can also be written in a form of triangle inequality

$$\sqrt{\Gamma(\Xi_b^- \rightarrow \Lambda_b \pi^-)} + \sqrt{\Gamma(\Xi_c^0 \rightarrow \Lambda_c \pi^-)} \geq \sqrt{0.9 \times 10^{10} \text{ s}^{-1}} . \quad (13)$$

Although at present it is not possible to evaluate in a reasonably model independent way the matrix element in eq.(5) for the ‘spectator’ decay amplitude, the estimate (13) shows that at least some of the discussed pion transitions should go at the level of 0.01 ps^{-1} . Also the numerical result for ΔA_S invites one more remark, related to the problem of the differences of lifetimes of heavy hadrons. Namely, we have seen here that the values of the matrix elements x and y extracted from the data on the lifetimes of charmed hyperons, translate, in terms of the pion transitions, into a ‘natural’ magnitude of the decay amplitude, which one would expect, based on the experience with the decays of the ordinary strange hyperons. On the other hand, these phenomenological numerical values of x and y are deemed to be substantially enhanced with respect to the simple estimates [11, 8], $y = -x = f_D^2 m_D / 12 \approx 0.006 \text{ GeV}^3$, based on non-relativistic-like ideas about the structure of the heavy baryons and mesons. The discussed here relation of these matrix elements to the decays $\Xi_Q \rightarrow \Lambda_Q \pi$ and the evaluation of their significance, perhaps, tell us that the phenomenological values of x and y [4] are of a ‘normal’ magnitude, while the early simplistic theoretical estimates were simply too low.

To summarize. It is shown that the strangeness-changing decays $\Xi_Q \rightarrow \Lambda_Q \pi$ of the b and c cascade hyperons should go at the rate of order 0.01 ps^{-1} and that they should obey the $\Delta I = 1/2$ rule. The decays $\Xi_b \rightarrow \Lambda_b \pi$ are purely S wave in the heavy quark limit and their amplitude, as well as the S wave amplitude of the decays $\Xi_s \rightarrow \Lambda_s \pi$, is expressed by the current algebra in terms of matrix elements of four-quark operators over the heavy baryons. The difference between the S wave amplitude of the charmed hyperon decay and that of the Ξ_b is related, within the heavy quark expansion, to the differences of lifetimes among the heavy hyperons. The semileptonic transitions $\Xi_Q \rightarrow \Lambda_Q \ell \nu$ are of a purely $0^+ \rightarrow 0^+$ type and thus have $g_A = 0$, $g_V = 1$, while the radiative weak decay $\Xi_b \rightarrow \Lambda_b \gamma$ is shown to be

greatly suppressed in the heavy quark limit.

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